Light Quark Energy Distribution in Heavy Quark Symmetry

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Abstract

Using heavy quark symmetry we construct an energy sum rule relating the form factor of a heavy flavored meson $\xi(w)$ to the light quark energy distribution. We find that the current available data for $\xi(w)$ is consistent with a broad energy distribution rather than with that of a pure quark/anti-quark bound state, indicating a large nonvalence content.

12.38Lg, 13.20Jf, 14.40Jz

There is growing interest in the physics of heavy quark systems stimulated by the recognition [1,2] of the important role of the heavy quark symmetry of the QCD lagrangian in the infinite quark mass limit. Interest has also been fostered by recent measurements of many heavy hadron decay modes [3] which provide more stringent challenges of our understanding of the interplay between strong and weak interactions [4]. In the heavy quark mass limit, $m \to \infty$, the spectrum of heavy mesons containing one heavy and one light valence quarks, become degenerated and independent of the polarization state of the light degrees of freedom [2]. Furthermore, an effective decoupling of the dynamics of a heavy quark reduces the problem to that of a cloud of light quarks and gluons in an external colored source [2,5,6]. These two facts imply definite relations between heavy meson matrix elements of various currents that involve heavy quark field operators. In particular it has been shown that in the limit of an exact heavy quark symmetry all form factors associated with such matrix elements are determined by a single, universal function $\xi(w)$, called the Isgur-Wise form factor [2,7]. The function $\xi(w)$ is usually introduced through a matrix element of a vector heavy quark current evaluated between two $J^P = 0^-$ heavy meson states,

$$\langle p'|\overline{\Psi}(0)\gamma^\mu\Psi(0)|p\rangle = f_+(t)(p'+p)^\mu + f_-(t)(p'-p)^\mu, \tag{1}$$

with $p'^2 = M'^2$, $p^2 = M^2$ and $t = (p'-p)^2$. In the $m, M, M' \to \infty$ limit it is useful to introduce the velocity transfer, $t^2/m^2 = (v'-v)^2$. Defining $w \equiv v \cdot v' = 1 - t/2m^2$ the form factors $f_\pm(t)$, which are not independent reduce to a single function of w,

$$\begin{aligned} f_+(t) &\to \frac{M'+M}{\sqrt{4M'M}}\xi(w) \to \xi(w), \\ f_-(t) &= \frac{M-M'}{M+M'}f_(t) \to 0. \end{aligned} \tag{2}$$

There have been many attempts to calculate the universal function $\xi(w)$ [6,8,9] and also the $O(1/m)$ corrections to $f_\pm$ [10,11]. As already pointed out, due to the decoupling of the heavy quark, $\xi(w)$ is primarily determined by the dynamics of light degrees of freedom thereby considerably simplifying the problem. This point is implicit in the constituent quark model calculations [8] in which $\xi(w)$ is determined by a noncovariant wave function of a light

spectator quark. Another approach is based on a QCD sum rule analysis of the vacuum three-point correlator [6]. The standard interpolation method is used to relate the spectral representation of the correlator to the contribution from the ground state, thus to $\xi(w)$. In such an approach the decoupling of the heavy quark makes $\xi(w)$ depend on the propagator of the light quark in an external field generated by the heavy quark.

In this paper we are interested in studying the structure of the distribution of light degrees of freedom in the heavy meson, and propose an approach in which we relate $\xi(w)$ to the nondiagonal correlator of two currents, i.e. with only one of the meson states in Eq. (1) being interpolated. The correlator is then evaluated between the vacuum and the other heavy meson state. It is in a sense an intermediate approach between the constituent quark model and a three-point function QCD sum rule calculation, and it will become clear that such an approach will allow us to relate, in a covariant way, the universal function $\xi(w)$ to gauge invariant matrix elements describing light quark and gluon content of a heavy meson. The general form of a correlator which we shall be dealing with in the rest of this paper is

$$\delta_{ij} T(p,q) = i \int dz e^{iqz} \langle p, i | T O(\frac{z}{2}) O_j(-\frac{z}{2}) | 0 \rangle, \tag{3}$$

where $|p,i\rangle$ is a heavy meson state with momentum p, ($p^2 = M^2$) containing a light valence quark with flavor i, and the operators O, O_i are given by $O(z) = \overline{\Psi}(z)\Gamma\Psi(z)$, $O_i(z) = \overline{\Psi}(z)\Gamma'\psi_i(z)$. Here Ψ and ψ_i are the heavy and light quark field operators and Γ and Γ' are some combination of Dirac gamma matrices, respectively. We start with the phenomenological analysis of the correlator $T(p,q)$. The 2-dimensional plane in variables $q_{1,2}^2$, ($q_{1,2} = p/2 \pm q$), is conventionally parameterized in terms of ω, λ and q^2, i.e. $q_{1,2}^2 = (1 \pm \omega) q^2 \pm \lambda + \frac{p^2}{4}$. For fixed ω and λ, with $|\omega| \leq 1$ and $\omega\lambda > p^2/4$, $T(p,q) = T(\omega, \lambda, q^2)$ satisfies a dispersion relation in the variable q^2 [12] with

$$\delta_{ij} Im T(\lambda, \omega, q^2) = \frac{1}{2} \int dz e^{iqz} \left[\langle p, i | O(\frac{z}{2}) O_j(-\frac{z}{2}) | 0 \rangle + \langle p, i | O_j(-\frac{z}{2}) O(\frac{z}{2}) | 0 \rangle \right]. \tag{4}$$

Inserting a complete set of intermediate states with masses M_n in Eq. (4) gives,

$$\delta_{ij} T(\lambda, \omega, q^2)$$

$$= \sum_n \frac{\theta\left(\frac{M_n^2-\lambda-p^2/4}{1+\omega} - s_{min}\right)}{M_n^2 - \lambda - p^2/4 - (1+\omega)q^2 - i\epsilon} \langle p,i|O(0)|n,q_{1n}\rangle\langle n,q_{1n}|O_j(0)|0\rangle_{q_{1n}^2=(p/2+q_n)^2=M_n^2}$$
$$+\frac{\theta\left(\frac{M_n^2+\lambda-p^2/4}{1-\omega} - s_{min}\right)}{M_n^2 + \lambda - p^2/4 - (1-\omega)q^2 - i\epsilon} \langle p,i|O_j(0)|n,q_{2n}\rangle\langle n,q_{2n}|O(0)|0\rangle_{q_{2n}^2=(p/2-q_n)^2=M_n^2}. \tag{5}$$

where $s_{min} = \min\left(\frac{-\lambda-p^2/4}{1+\omega}, \frac{\lambda-p^2/4}{1-\omega}\right)$. With our choice of the quark content for the operators O, O_j the first term in the RHS of Eq. (5) has contributions from resonances containing one heavy and one light valence quark, while the second term involves states with two heavy quarks and does not contribute to $\xi(w)$. Therefore, we shall choose the variables ω, λ, q^2 in such a way that the second term is nonleading in the $1/m$ expansion. This corresponds to the choice $\lambda \gtrsim p^2/4$, $(\omega > 0)$ and $q^2 \sim p^2/2(1+\omega)$ which leads to a suppression by a factor of $O(1/m)$ of the second term with respect to the first term in Eq. (5). Furthermore it can be shown that the $(1-\omega)$ dependence is also nonleading so that in the leading order analysis we set $\lambda \gtrsim p^2/4$ and $\omega = 1$. It is also more convenient to use, in the infinite heavy quark mass limit [13,14], the spectral representation in the energy variables instead of the invariant mass. With the resonance fixed at $q_{1n}^2 = M_n^2$ the momentum transfer is $t = 2m^2(1-w) = (p-q_{1n})^2 = -\lambda + p^2/4$ and the correlator will be evaluated for $q_1^2 = (m+E_q)^2 \sim m^2 + 2mE_q$. Denoting the energy of an excited state by E_n, $M_n^2 \sim m^2 + 2mE_n$, the correlator in the variables E_q and w for $\omega = 1$ and $\lambda \gtrsim p^2/4$ is given by

$$\delta_{ij} T(w, E_q) = \sum_n \frac{\langle p,i|O(0)|n,q_{1n}\rangle\langle n,q_{1n}|O_j(0)|0\rangle_{q_{1n}^2=M_n^2,(p-q_{1n})^2=2m^2(1-w)}}{2m(E_n - E_q - i\epsilon)}. \tag{6}$$

In the following we restrict our analysis to two choices of the operators O, O_j given by two sets of the Γ, Γ' matrices; Set 1 : $\Gamma = \gamma^\mu,\ \Gamma' = \gamma_5$, Set 2 : $\Gamma = 1,\ \Gamma' = \gamma_5$. We will explicitly show that to leading order in the $1/m$ expansion the two sets lead to a unique sum rule for the form factor $\xi(w)$. For $|E_q - E_n| >> 0$ a standard assumption is that the contributions to the sum in Eq. (6) from the ground state and from higher resonances can be parameterized by a single state with energy E_0 and zero width together with a continuum starting at

with threshold energy E_c For two Γ sets the leading contribution to the phenomenological correlator as $m \to \infty$ is given by

$$
\begin{aligned}
\text{Set 1}: T^\mu(w, E_q) &= \frac{i}{2}(\frac{3}{2}p + q)^\mu \left[\frac{f_h \xi(w)}{(E_q - E_0 + i\epsilon)} + T_c(\tilde{t}, E_q, E_c)\right], \\
\text{Set 2}: T(w, E_q) &= \frac{i}{2} m \left[\frac{f_h (1+w) \xi(w)}{(E_q - E_0 + i\epsilon)} + T_c(\tilde{t}, E_q, E_c)\right],
\end{aligned} \tag{7}
$$

where f_h is a ground state heavy meson decay constant, E_c is the effective continuum threshold energy, and T_c is the contribution from the continuum whose explicit form will be discussed later. In deriving the second equation above we have used $\langle p' | \overline{\Psi}(0) \Psi(0) | p \rangle = m(1+w)\xi(w)$, which follows from Eq. (1) taking the derivative of the vector current in the $M' \to M$ limit.

We shall now discuss the theoretical description of the correlators. For $E_q >> E_0 \sim \Lambda_{QCD}$ the heavy quark at the intermediate state in the correlator in Eq. (3) is off its energy shell roughly by E_q and the singularity of the perturbative heavy quark propagator dominates the spacetime behavior of the operator product in Eq. (3). In a fixed gauge, the leading perturbative contribution is given by [18]

$$
\langle p, i | T \overline{\Psi}(\frac{z}{2}) \Gamma \Psi(\frac{z}{2}) \overline{\Psi}(-\frac{z}{2}) \Gamma' \psi_j(-\frac{z}{2}) | 0 \rangle \to \delta_{ij} \mathrm{Tr} \left[\langle p, i | \overline{\Psi}(\frac{z}{2}) \psi_i(-\frac{z}{2}) | 0 \rangle \Gamma S(z) \Gamma' \right], \tag{8}
$$

where $S(z)$ is the perturbative heavy quark propagator and the trace is taken over the spinor indices. The nonperturbative part of the propagator leads to a $O(1/m)$ correction and is not taken into account in the leading order analysis. For the two sets of operators discussed earlier the correlators are given by

$$
\begin{aligned}
T^\mu(p, q) &= i f_h \int d^4 z \int \frac{d^4 k}{(2\pi)^4} \frac{e^{-i(q-k)z}}{k^2 - m^2 + i\epsilon} \Big[(k^\mu + p^\mu) m \phi(z^2, z \cdot p) \\
&\qquad + k^\mu m \phi_-(z^2, z \cdot p) + k^\nu m \phi_{T\nu\mu}(z^2, z \cdot p) \Big], \\
T(p, q) &= i f_h \int d^4 z \int \frac{d^4 k}{(2\pi)^4} \frac{e^{-i(q-k)z}}{k^2 - m^2 + i\epsilon} \left[(k \cdot p + m^2) \phi(z^2, z \cdot p) + m^2 \phi_-(z^2, z \cdot p) \right],
\end{aligned} \tag{9}
$$

where

$$\phi^\dagger(z^2, z\cdot p) \equiv \frac{-i}{f_h}\langle 0|\overline{\psi}_i(\frac{z}{2})\frac{\not p}{p^2}\gamma_5\Psi(-\frac{z}{2})|p,i\rangle,$$
$$\phi^\dagger_-(z^2, z\cdot p) \equiv \frac{i}{f_h}\left[\frac{1}{m}\langle 0|\overline{\psi}_i(\frac{z}{2})\gamma_5\Psi(-\frac{z}{2})|p,i\rangle + \langle 0|\overline{\psi}_i(\frac{z}{2})\frac{\not p}{p^2}\gamma_5\Psi(-\frac{z}{2})|p,i\rangle\right],$$
$$\phi^\dagger_{T\nu\mu}(z^2, z\cdot p) \equiv \frac{-i}{f_h m}\langle 0|\overline{\psi}_i(\frac{z}{2})i\sigma_{\nu\mu}\gamma_5\Psi(-\frac{z}{2})|p,i\rangle, \tag{10}$$

and the normalization of the leading amplitude, $\phi(z^2, z\cdot p)$, is given by $\phi(z^2, z\cdot p)_{z=0} = 1$. In the infinite heavy quark mass limit the operators that depend on a spin structure of the light quark (i.e, $\overline{\Psi}\gamma^\mu\gamma_5\psi_i$ and $\overline{\Psi}\gamma_5\psi_i$) are degenerate [7] and thus lead to the same matrix element. For this reason ϕ_- does not contribute in the leading order. Similarly it is seen that ϕ_T is suppressed by a power of $1/m$ relative to ϕ so only terms proportional to ϕ will be retained in the leading order analysis. The matrix element which defines ϕ can be rewritten as

$$\phi^\dagger(z^2, z\cdot p) = \frac{1}{if_h}e^{i\frac{p}{2}\cdot z}\langle 0|\overline{\psi}_i(z)\frac{\not p}{p^2}\gamma_5\Psi(0)|p,i\rangle \equiv e^{i\frac{p}{2}\cdot z}\phi_q^\dagger(z_T^2, z_L) \tag{11}$$

with the subscript L indicating the magnitude of a 4-vector projection parallel to p^μ and subscript T indicating the remaining three components of a 4-vector perpendicular to p^μ. The structure of the light quark wave function, ϕ_q is revealed by expanding in a Taylor series at $z = 0$,

$$\phi_q(z_T^2, z_L) = \sum_n \frac{(iz_L\sqrt{p^2})^n}{n!}\left[\langle\phi_q^n\rangle(0) + z_T^2\langle\phi_q^n\rangle'(0) + \frac{z_T^4}{2}\langle\phi_q^n\rangle''(0) + \cdots\right]. \tag{12}$$

The moments $\langle\phi_q^n\rangle^{(m)}(0)$ can be related to local matrix elements with n longitudinal, D_L, and $2m$ transverse, D_T, covariant derivative insertions, $D^\mu = \partial^\mu + igA^\mu(z)$ [14]. If we imagine that the heavy meson is composed of a single light quark bound by a chromoelectric potential of a static heavy quark, i.e. with no dynamical gluons, then $D_L = \partial_L + igA_L(z_T)$ being z_L independent implies that $\langle\phi_q^n\rangle^{(m)}(0)$ with fixed m determine the energy (longitudinal component of the 4-vector) distribution of a single light quark. The heavy meson having a definite energy, E_0, requires the light quark to be on the energy shell as well, and as a consequence $\langle\phi_q^n\rangle(0) = (E_0/\sqrt{p^2})^n$ or

$$\phi_q(z_T^2, z_L) = e^{iz_L E_0}\phi_q(z_T^2). \tag{13}$$

The transverse amplitude $\phi_q(z_T^2)$ is the bound state wave function which has a nontrivial behavior for $|z_T| \sim 1/E_0$ even in the static case and it approaches a plane wave solution only in a free case i.e. when $A(z_T) \to 0$. However the z_L dependent amplitude is the same whenever ϕ_q describes a bound state or a free particle. Since in general the vector potential $A = A(z_L, z_T)$ is nonstatic one expects a sizable contributions from the gluon and/or $q\overline{q}$ sea. The z_L dependence of A_L implies that the spectral representation of $\phi_q(z_L)$ is a smeared distribution around $E = E_0$ rather then being proportional to a delta function. Furthermore since the gluon field $A^\mu = A^\mu(z)$ should be in general p^μ independent, the division between the longitudinal and transverse wave functions is purely a matter of convenience. Also relevant is the scale governing the energy and momentum smearing intervals of the two wave functions which should be of the same order. Thus, despite the fact that in leading order the correlators in Eq. (9) are not sensitive to the transverse wave function, the above argument still permits conclusions on the size of the gluon momentum distribution from the analysis of the longitudinal wave function alone.

Transforming to the momentum representation the correlators are given by

$$\begin{aligned} T^\mu(w, E_q) &= \frac{i}{2} f_h (\frac{3}{2}p + q)^\mu \int_0 \frac{dE}{2\pi} \frac{\phi_q(E)}{E_q - Ew + i\epsilon}, \\ T(w, E_q) &= \frac{i}{2} f_h m (1 + w) \int_0 \frac{dE}{2\pi} \frac{\phi_q(E)}{E_q - Ew + i\epsilon} \end{aligned} \tag{14}$$

with $\phi_q(E) = \int dz_L e^{iEz_L} \phi(z_T^2 = 0, z_L)$ and $\int_0 \frac{dE}{2\pi} \phi_q(E) = 1$. The leading order theoretical correlators of Eq. (14) are also used to define a continuum contribution T_c introduced in Eq. (7). The standard assumption is that the spectral density of $T_c(\tilde{t}, E_q, E_c)$ representing the sum over higher resonances can be replaced by the theoretical one at $E > E_c$. From Eqs. (7) and (14) we obtain a single sum rule for $\xi(w)$.

$$\int_0 \frac{dE}{2\pi} \frac{\phi_q(E)}{E_q - Ew + i\epsilon} = \frac{\xi(w)}{E_q - E_0 + i\epsilon} + \int_{\frac{E_c}{w}} \frac{dE}{2\pi} \frac{\phi_q(E)}{E_q - Ew + i\epsilon}. \tag{15}$$

In order to reduce the contributions from the unknown higher order terms in the expansion in Eq. (12) and suppress the sensitivity of the sum rule to the phenomenological parameterization of the spectral density, a standard Borel transformation ($E_q \to T$) is performed for

Eq. (15) leading to

$$
\begin{aligned}
& e^{-wE_0/T} + \int_0 \frac{dE}{2\pi} e^{-wE/T} \left[\phi_q(E) - 2\pi\delta(E - E_0)\right] \\
& = \xi(w) e^{-E_0/T} + \int_{\frac{E_c}{w}} \frac{dE}{2\pi} e^{-wE/T} \phi_q(E). \qquad (16)
\end{aligned}
$$

As explained earlier in a simple quantum mechanical description in which a heavy meson contains a single quark orbiting around a static chromoelectric source, $\phi_q(E) = 2\pi\delta(E - E_0)$ and since $E_c > E_0$, the sum rule automatically leads to a correct normalization, $\xi(1) = 1$, for all values of T. At first look the above result may seem dependent on the choice of the parameterization of the phenomenological spectral density. This is however not the case. If a more detailed parameterization were used, involving explicitly higher resonances instead of a smooth continuum approximation, due to orthogonality of the mass eigenstates such resonances would not contribute to the RHS of Eq. (16) at $w = 1$ leaving $\xi(1) = 1$ coming only from the ground state contribution. In a realistic situation $\phi_q(E)$ is a smeared distribution around $E = E_0$. Furthermore it follows from the analysis of the normalization, $f_{H \to L}$, of the heavy-to-light meson matrix elements that the $f_{H \to L} \sim O(m^{-3/2})$ behavior [15] requires $\phi_q(E)$ to vanish only linearly with E for $E/E_0 << 1$ [16]. In order to be consistent with the phenomenological parameterization of the spectral density we shall assume that $\phi_q(E)$ vanishes for $E > E_c$. The value of the threshold energy has been obtained in Ref. [6,14] and it varies in the range $2E_0 \lesssim E_c \lesssim 3E_0$. We then use the following parameterization $\phi_q \propto E(E_c - E)\exp(-(E - E_0)^2/E_W^2)$ and study the form factor $\xi(w)$ for different values of the width parameter E_W. For a given value of E_c/E_0 and E_W/E_0 the Borel parameter (T/E_0) is fixed by the normalization $\xi(1) = 1$. In Fig. 1 we plot our predictions for $\xi(w)$ for $E_c/E_0 = 2.1, 2.5, 3$ given by the set of dashed, solid and dotted curves respectively and for $E_W/E_0 = .1$ and 10 corresponding to upper and lower curves in each set (dashed, solid, dotted) respectively. As E_W decreases the predictions are sensitive to both the detailed shape of the energy distribution of the light quark and to the value of the continuum threshold energy. On the other hand, for increasing E_W our predictions become weakly dependent on E_c/E_0 and E_W/E_0, and are bound from below by $1/w^2$ behavior which in turn follows

from the linear behavior of $\phi_q(E)$ at small E. The data corresponding to the $B \to D$ form factor (finite m) and extracted from the $B \to Dl\overline{\nu}_l$ with $A = (\tau_B/1.18) * (|V_{cb}|/0.05) = 1.11$ have been taken from Ref. [19]. The comparison of our results with the experimental data shows that the $\xi(w)$ form factor is inconsistent with the usually assumed peaking (δ-type) approximation to the distribution amplitude corresponding to small E_W/E_0) [20] and suggests a rather broad, $E_W > E_0$, energy distribution. This in turn implies large gluon and possibly sea quark amplitudes in a heavy meson and we conclude that a large fraction of the energy-momentum of the light degrees of freedom is distributed over the nonvalence components. Although our results rely on a comparison with heavy, but finite, quark mass data, both the $O(1/m)$ and perturbative, $O(\alpha(E))$ corrections, as shown in Ref. [11], are expected to be very small not exceeding a few percent.

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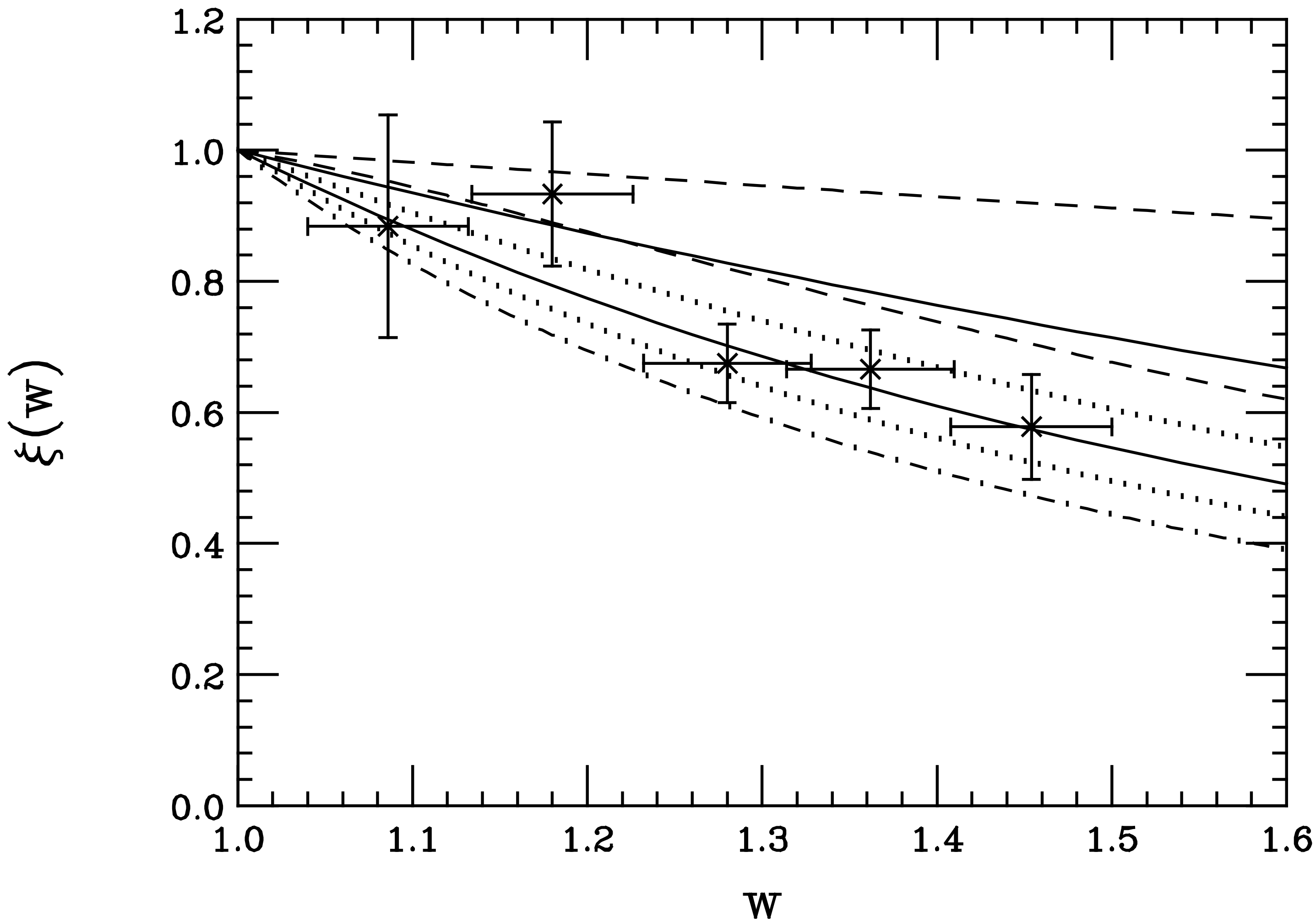
FIGURES

FIG. 1. The Isgur-Wise form factor $\xi(w)$. The dash-doted line is the $1/w^2$ asymptotic behavior. The variuos theoretical curves are explainded in the text.


ξ(w)
w
1.2
1.0
0.8
0.6
0.4
0.2
0.0
1.0
1.1
1.2
1.3
1.4
1.5
1.6